

The πNN system — recent progress

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Recent progress towards an understanding of the πNN system within chiral perturbation theory is reported. The focus lies on an effective field theory calculation and its comparison to phenomenological calculations for the reaction $NN \rightarrow d\pi$. In addition, the resulting absorptive and dispersive corrections to the πd scattering length are discussed briefly.

Keywords: πd scattering, pion production

1. Introduction

Pion reactions on few-nucleon systems provide access to various physics phenomena: deuterons can be used as effective neutron targets, null-experiments for isospin violation can be designed, and they are an important test of our understanding of the nuclear structure. What is therefore necessary is a controlled theoretical framework — a proper effective field theory needs to be constructed.

A first step in this direction was taken by Weinberg already in 1992.¹ He suggested that all that needs to be done is to convolute transition operators, calculated perturbatively in standard chiral perturbation theory (ChPT), with proper nuclear wave functions to account for the non-perturbative character of the few-nucleon systems. This procedure combines the distorted wave born approximation, used routinely in phenomenological calculations, with a systematic power counting for the production operators. Within ChPT this idea was already applied to a large number of reactions like $\pi d \rightarrow \pi d$,² $\gamma d \rightarrow \pi^0 d$,^{3,4} $\pi^3\text{He} \rightarrow \pi^3\text{He}$,⁵ $\pi^- d \rightarrow \gamma nn$,⁶ and $\gamma d \rightarrow \pi^+ nn$,⁷ where only the most recent references are given.

The central concept to be used in the construction of the transition operators is that of reducibility, for it allows one to disentangle effects of

the wave functions and those from the transition operators. As long as the operators are energy independent, the scheme can be applied straight forwardly,⁸ however, as we will see below, for energy dependent interactions more care is necessary.

Using standard ChPT especially means to treat the nucleon as a heavy field. Corrections due to the finite nucleon mass, M_N , appear as contact interactions on the lagrangian level that are necessarily analytic in M_N . However, some pion–few-nucleon diagrams employ few–body singularities that lead to contributions non–analytic in m_π/M_N , with m_π for the pion mass. In Ref.⁹ it is explained how to deal with those.

A problem was observed when the original scheme by Weinberg was applied to the reactions $NN \rightarrow NN\pi$:^{10,11} Potentially higher order corrections turned out to be large and lead to even larger disagreement between theory and experiment. For the reaction $pp \rightarrow pp\pi^0$ one loop diagrams that in the Weinberg counting appear only at NNLO were evaluated^{12,13} and they turned out to give even larger corrections putting into question the convergence of the whole series. However, already quite early the authors of Refs.^{14,15} stressed that an additional new scale enters, when looking at reactions of the type $NN \rightarrow NN\pi$, that needs to be accounted for in the power counting. Since the two nucleons in the initial state need to have sufficiently high kinetic energy to put the pion in the final state on-shell, the initial momentum needs to be larger than

$$p_{thr} = \sqrt{M_N m_\pi} .$$

The proper way to include this scale was presented in Ref.¹⁶ and implemented in Ref.¹⁷ — for a recent review see Ref.¹⁸ As a result, pion p -waves are given by tree level diagrams up to NNLO in the modified power counting and the corresponding calculations showed satisfying agreement with the data.¹⁶ However, for pion s -waves loops appear already at NLO. In the next section we will discuss their effect on the reaction $NN \rightarrow d\pi$ near threshold. In some detail we will compare the effective field theory result to that of phenomenological calculations.

Since the Delta–nucleon mass difference, Δ , is numerically of the order of p_{thr} , also the Delta–isobar should be taken into account explicitly as a dynamical degree of freedom.¹⁴ We will use a scheme where

$$\Delta \sim p_{thr} .$$

Once the reaction $NN \rightarrow d\pi$ is understood within effective field theory one is in the position to also calculate the so-called dispersive and

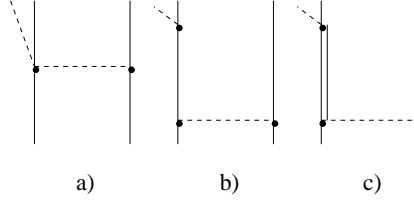


Fig. 1. Tree level diagrams that contribute to $pp \rightarrow d\pi^+$ s -waves up to NLO. Solid lines denote nucleons, dashed ones pions and the double line the propagation of a Delta-isobar.

absorptive corrections to the πd scattering length. This calculation will be presented in section 3.

We close with a brief summary and outlook.

2. $NN \rightarrow d\pi$

The tree level amplitudes that contribute to $pp \rightarrow d\pi^+$ are shown in Fig. 1. In Ref.¹⁷ all NLO contributions of loops that start to contribute^a to $NN \rightarrow NN\pi$ at NLO were calculated in threshold kinematics — that is neglecting the distortions from the NN final- and initial state interaction and putting all final states at rest. At threshold only two amplitudes contribute, namely the one with the nucleon pair in the final and initial state in isospin 1 (measured, e.g., in $pp \rightarrow pp\pi^0$) and the one where the total NN isospin is changed from 1 to 0 (measured, e.g., in $pp \rightarrow d\pi^+$)^b. It was found that the sum of all loops that contain Delta-excitations vanish in both channels. This was understood, since the loops were divergent and at NLO no counter term is allowed by chiral symmetry. On the other hand the nucleonic loops were individually finite. It was found that the sum of all nucleonic loops that contribute to $pp \rightarrow pp\pi^0$ vanish, whereas the sum of those that contribute to $pp \rightarrow d\pi^+$ gives a finite answer. The resulting amplitude grows linear with the initial momentum.

At that time it appeared as a puzzle why the loops vanished for the reaction $pp \rightarrow pp\pi^0$ — no obvious symmetry reason could be identified. However, in Ref.¹⁹ it was pointed out that the linear growth of the amplitude

^aIn a scheme with two expansion parameters — here m_π and p_{thr} — loops no longer contribute at a single order but in addition to all orders higher than where they start to contribute.

^bThe third independent amplitude, where the NN isospin is changed from 0 to 1 in the production process and that can be extracted from $pn \rightarrow pp\pi^-$, vanishes at threshold as a consequence of selection rules.

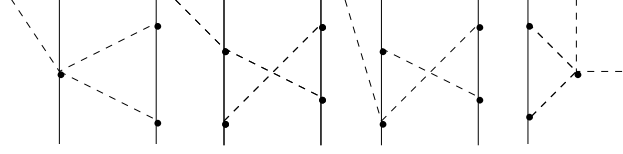


Fig. 2. Irreducible pion loops with nucleons only that start to contribute to $NN \rightarrow NN\pi$ at NLO that were considered in Ref.¹⁷

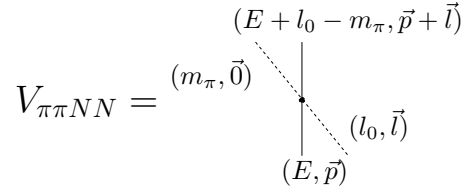


Fig. 3. The $\pi N \rightarrow \pi N$ transition vertex: definition of kinematic variables as used in the text.

for the charged pion production is the problematic one: when evaluated for finite outgoing NN momenta, the transition amplitudes turned out to scale as the momentum transfer. Especially, the amplitudes then grow linearly with the external NN momenta. As a consequence, once convoluted with the NN wavefunctions, a large sensitivity to those was found, in conflict with general requirements from field theory. In light of these insights it was acknowledged that the loops for $pp \rightarrow d\pi^+$ were the ones not understood. The solution to this puzzle was presented in Ref.²⁰ and will be reported now.

The observation central to the analysis is that the leading $\pi N \rightarrow \pi N$ transition vertex, as it appears in Fig. 1a, is energy dependent. Using the notation of Fig. 3 its momentum and energy dependent part may be written as^c

$$\begin{aligned}
 V_{\pi\pi NN} &= l_0 + m_\pi - \frac{\vec{l} \cdot (2\vec{p} + \vec{l})}{2M_N} \\
 &= \underbrace{2m_\pi}_{\text{on-shell}} + \underbrace{\left(l_0 - m_\pi + E - \frac{(\vec{l} + \vec{p})^2}{2M_N} \right)}_{(E' - H_0) = (S')^{-1}} - \underbrace{\left(E - \frac{\vec{p}^2}{2M_N} \right)}_{(E - H_0) = S^{-1}}. \quad (1)
 \end{aligned}$$

^cThe expressions for the vertices can be found in Ref.²¹ Note that the $\pi N \rightarrow \pi N$ vertex from $\mathcal{L}_{\pi N}^{(1)}$ as well as its recoil correction from $\mathcal{L}_{\pi N}^{(2)}$ are to be used already at leading order as a consequence of the modified power counting.

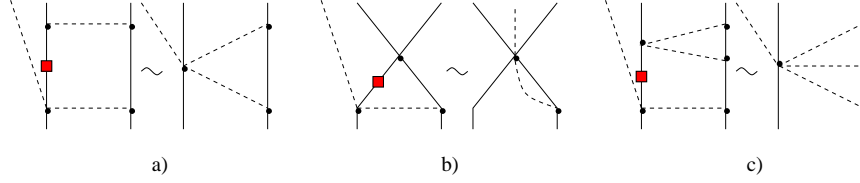


Fig. 4. Induced irreducible topologies, when the off-shell terms of Eq. (1) hit the NN potential in the final state. The filled box on the nucleon line denotes the propagator canceled by the off-shell part of the vertex.

For simplicity we skipped the isospin part of the amplitude. The first term in the last line denotes the transition in on-shell kinematics, the second the inverse of the outgoing nucleon propagator and third the inverse of the incoming nucleon propagator. First of all we observe that for on-shell incoming and outgoing nucleons, the $\pi N \rightarrow \pi N$ transition vertex takes its on-shell value $2m_\pi$ — even if the incoming pion is off-shell, as it is for diagram *a* of Fig. 1. This is in contrast to standard phenomenological treatments,²² where l_0 was identified with $m_\pi/2$ — the energy transfer in on-shell kinematics — and the recoil terms were not considered. Note, since $p_{thr}^2/M_N = m_\pi$ the recoil terms are to be kept.

The second consequence of Eq. (1) is even more interesting: when the $\pi N \rightarrow \pi N$ vertex gets convoluted with NN wave functions, only the first term leads to a reducible diagram. The second and third term, however, lead to irreducible contributions, since one of the nucleon propagators gets canceled. This is illustrated in Fig. 4, where those induced topologies are shown that appear, when one of the nucleon propagators is canceled (marked by the filled box) in the convolution of typical diagrams of the NN potential with the $NN \rightarrow NN\pi$ transition operator. Power counting gives that diagrams *b* and *c* appear only at order $N^4\text{LO}$ and $N^3\text{LO}$, respectively. However, diagram *a* starts to contribute at NLO and it was found in Ref.²⁰ that those induced irreducible contributions cancel the finite remainder of the NLO loops in the $pp \rightarrow d\pi^+$ channel. Thus, up to NLO only the diagrams of Fig. 1 contribute to $pp \rightarrow d\pi^+$, with the rule that the $\pi N \rightarrow \pi N$ vertex is put on-shell.

The result found in Ref.²⁰ is shown in Fig. 5, where the total cross section (divided by the energy dependence of phase space) is plotted against the normalized pion momentum. The dashed line is the result of the model by Koltun and Reitan,²² as described above, whereas the solid line shows the result of the ChPT calculation of Ref.²⁰.

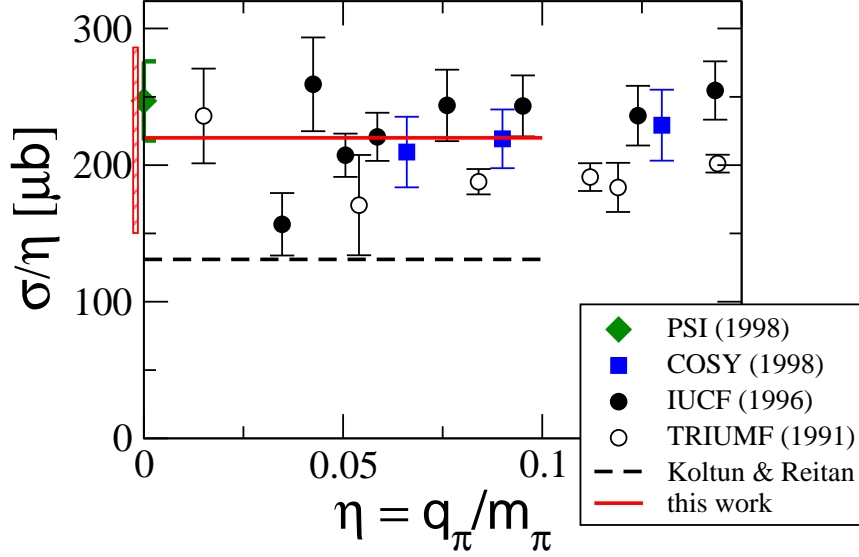


Fig. 5. Comparison of our results to experimental data for $NN \rightarrow d\pi$. The dashed line corresponds to the model of Koltun and Reitan,²² whereas the solid line is the result of the ChPT calculation of Ref.²⁰ The estimated theoretical uncertainty (see text) is illustrated by the narrow box. The data is from Refs.²⁵ (open circles),²⁶ (filled circles) and²⁷ (filled squares). The first data set shows twice the cross section for $pn \rightarrow d\pi^0$ and the other two the cross section for $pp \rightarrow d\pi^+$.

3. Comparison to phenomenological works

Phenomenological calculations for the reaction $pp \rightarrow d\pi^+$ in near threshold kinematics are given, e.g., in Ref.²³ and Ref.²⁴. In both works in addition to the diagrams of Ref.²² some Delta-loops as well as short range contributions are included — heavy meson exchanges for the former and off-shell πN scattering^d for the latter. Based on this the cross section for $pp \rightarrow d\pi^+$ is now even overestimated near threshold. How can we interpret this discrepancy in light of the discussion above?

First of all, the NLO parts of the Delta-loops cancel, as was shown already in Ref.¹⁷. However, in both Refs.^{23,24} only one of these diagrams was included and, especially for Ref.²³ gave a significant contribution. The only diagram of those NLO loops shown in Fig. 2 that is effectively included in Ref.²⁴ is the fourth, since the pion loop there can be regarded as part of the $\pi N \rightarrow \pi N$ transition T -matrix. However, as described, the contri-

^dThat those are also short range contributions is discussed in Ref.¹⁸

bution of this diagram gets canceled by the others shown in Fig. 2 and the induced irreducible pieces described above. Therefore, the physics that enhances the cross section compared to the work of Ref.²² in Refs.^{23,24} is completely different to that of Ref.²⁰ — the phenomenological calculations miss the essential contribution and are in conflict with both field theoretic consistency and chiral symmetry.

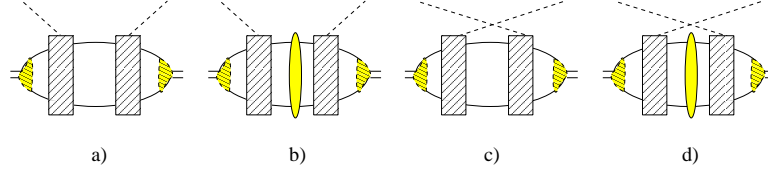
What are the observable consequences of the difference between the ChPT calculation and the phenomenological ones? As explained, in the former the near threshold cross section for $pp \rightarrow d\pi^+$ is basically given by a long-ranged pion exchange diagram, whereas the latter rely on short ranged operators with respect to the NN system. Obviously those observables are sensitive to this difference that get prominent contributions from higher partial waves in the final NN system. We therefore need to look at the reaction $pp \rightarrow pn\pi^+$. Unfortunately, the total cross section for this reaction is largely saturated by NN S -waves in the final state (see, e.g., Fig. 17 in Ref.¹⁸). On the other hand, linear combinations of double polarization observables allow one to remove the prominent components and the sub-leading amplitudes should be visible. We therefore expect from the above considerations that the phenomenological calculations give good results for polarization observables for $pp \rightarrow d\pi^+$, whereas there should be deviations for some of those for $pp \rightarrow pn\pi^+$. Predictions for these observables were presented in Ref.²⁸ and indeed the π^+ observables with the deuteron in the final state are described well whereas there are discrepancies for the pn final state (see Fig. 24 of Ref.¹⁸).

It remains to be seen how well the same data can be described in the effective field theory framework. Up to NNLO the number of counter terms is quite low: there are two counter terms for pion s -waves, that can be arranged to contribute to $pp \rightarrow pp\pi^0$ and $pp \rightarrow d\pi^+$ individually, and then there is one counter term for pion p -waves, that contributes only to a small amplitude in charged pion production.¹⁶ On the other hand there is a huge amount of even double polarized data available^{29–31} — and there is more to come especially for $pn \rightarrow pp\pi^-$.³²

4. Corrections to $a_{\pi d}$

The πd scattering length is known to a high accuracy from measurements on pionic deuterium³³

$$a_{\pi d}^{\text{exp}} = (-26.1 \pm 0.5 + i(6.3 \pm 0.7)) \times 10^{-3} m_{\pi}^{-1}, \quad (2)$$

Fig. 6. Dispersive corrections to the πd scattering length.

where m_π denotes the mass of the charged pion. In the near future a new measurement with a projected total uncertainty of 0.5% for the real part and 4% for the imaginary part of the scattering length will be performed at PSI.³⁴ What is striking with this result is the quite large imaginary part that may be written as

$$4\pi\text{Im}(a_{\pi d}) = \lim_{q \rightarrow 0} q \{ \sigma(\pi d \rightarrow NN) + \sigma(\pi d \rightarrow \gamma NN) \} , \quad (3)$$

where q denotes the relative momentum of the initial πd pair. The ratio $R = \lim_{q \rightarrow 0} (\sigma(\pi d \rightarrow NN)/\sigma(\pi d \rightarrow \gamma NN))$ was measured to be 2.83 ± 0.04 .³⁵ At low energies diagrams that lead to a sizable imaginary part of some amplitude are expected to also contribute significantly to its real part. Those contributions are called dispersive corrections. As a first estimate Brückner speculated that the real and imaginary part of these contributions should be of the same order of magnitude.³⁶ This expectation was confirmed within Faddeev calculations in Refs.³⁷ Given the high accuracy of the measurement and the size of the imaginary part of the scattering length, another critical look at this result is called for as already stressed in Refs.^{38,39}. A consistent calculation is only possible within a well defined effective field theory — the first calculation of this kind was presented in Ref.⁴⁰ and is briefly sketched here.

To identify the diagrams that are to contribute we first need to specify what we mean by a dispersive correction. We define dispersive corrections as contributions from diagrams with an intermediate state that contains only nucleons, photons and at most real pions. Therefore, all the diagrams shown in Fig. 6 are included in our work. On the other hand, all diagrams that, e.g., have Delta excitations in the intermediate state do not qualify as dispersive corrections, although they might give significant contributions.⁴¹

The hatched blocks in the diagrams of Fig. 6 refer to the relevant transition operators for the reaction $NN \rightarrow NN\pi$ depicted in Fig. 1. Also in the kinematics of relevance here the $\pi N \rightarrow \pi N$ transitions are to be taken with their on-shell value $2m_\pi$. Using the CD-Bonn potential⁴² for the NN

distortions we found for the dispersive correction from the purely hadronic transition

$$a_{\pi d}^{disp} = (-6.3 + 2 + 3.1 - 0.4) \times 10^{-3} m_{\pi}^{-1} = -1.6 \times 10^{-3} m_{\pi}^{-1}, \quad (4)$$

where the numbers in the first bracket are the individual results for the diagrams shown in Fig. 6, in order. There are two points important to stress, first of all the inclusion of the intermediate NN interaction is necessary (and required based on power-counting) and the crossed diagrams (diagram c and d) give a numerically significant contribution. The latter finding might come as a surprise on the first glance, however, please recall that in the chiral limit all four diagrams of Fig. 6 are kinematically identical and chiral perturbation theory is a systematic expansion around exactly this point. Thus, as a result we find that the dispersive corrections to the πd scattering length are of the order of 6 % of the real part of the scattering length. Note that the same calculation gave very nice agreement for the corresponding imaginary part.⁴⁰

In Ref.⁴⁰ also the electro-magnetic contribution to the dispersive correction was calculated. It turned out that the contribution to the real part was tiny — $-0.1 \times 10^{-3} m_{\pi}^{-1}$ — while the sizable experimental value for the imaginary part (c.f. Eqs. (2) and 3) was described well.

To get a reliable estimate of the uncertainty of the calculation just presented a NNLO calculation is necessary. At that order two counter terms appear for pions at rest that can be fixed from $NN \rightarrow NN\pi$, as indicated above. For now we need to do a conservative estimate for the uncertainty by using the uncertainty of order $2m_{\pi}/M_N$ one has for, e.g., the sum of all direct diagrams to derive a $\Delta a_{\pi d}^{disp}$ of around $1.4 \times 10^{-3} m_{\pi}^{-1}$, which corresponds to about 6% of $\text{Re}(a_{\pi d}^{\text{exp}})$. However, given that the operators that contribute to both direct and crossed diagrams are almost the same and that part of the mentioned cancellations is a direct consequence of kinematics, this number for $\Delta a_{\pi d}^{disp}$ is probably too large.

In Ref.⁴⁰ a detailed comparison to previous works is given. Differences in the values found for the dispersive corrections were traced to the incomplete sets of diagrams included.

5. Summary and Outlook

The process $NN \rightarrow NN\pi$ is a puzzle already since more than a decade. Given the progress presented above we have now reason to believe that this puzzle will be solved soon. This mentioned results could only be found, because a consistent effective field theory was used. For example, the potential

problem with the transition operators of Ref.,¹⁷ pointed at in Ref.,¹⁹ would always be hidden in phenomenological calculations, since the form factors routinely used there always lead to finite, well behaved amplitudes. The very large number of observables available for the reactions $NN \rightarrow NN\pi$ will provide a non-trivial test to the approach described.

Once the scheme is established, the same field theory can be used to analyze the isospin violating observables measured in $pn \rightarrow d\pi^{0,43}$ and $dd \rightarrow \alpha\pi^0$.⁴⁴ First steps in this direction were already done in Ref.⁴⁵ for the former and in Refs.^{46,47} for the latter.

Based on the calculation for $pp \rightarrow d\pi^+$ we also performed a calculation for the dispersive and absorptive corrections to the πd scattering length that were calculated for the first time within ChPT. The final answer turned out to be relatively small as a consequence of cancellations amongst various terms. This work is an important step forward towards a high accuracy calculation for the πd scattering length that will eventually allow for a reliable extraction of the isoscalar scattering length. However, before this can be done, isospin violating corrections⁴⁸ as well as the contributions from the Delta-isobar need to be evaluated.

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